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Cosmic-ray neutrinos from the decay of long-lived particle and the recent IceCube result



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ABSTRACT

Motivated by the recent IceCube result, we study high energy cosmic-ray neutrino flux from the decay of a long-lived particle. Because neutrinos are so transparent, high energy neutrinos produced in the past may also contribute to the present neutrino flux. We point out that the PeV neutrino events observed by IceCube may originate in the decay of a particle much heavier than PeV if its lifetime is shorter than the present cosmic time. It is shown that the mass of the particle responsible for the IceCube event can be as large as $\sim 10^{10}$ GeV. We also discuss several possibilities to acquire information about the lifetime of the long-lived particle.

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Neutrino astronomy provides a new window to the early universe. This is particularly because, contrary to particles with electromagnetic interactions (i.e., e^\pm , p and \bar{p} , γ and so on), neutrinos have very weak interaction and are very transparent. Thus, the cosmic-ray neutrino spectrum in the present universe contains various information about the production processes of energetic neutrinos.

Recently, the IceCube experiment reported the result of their analysis on high-energy neutrino events [1]. The IceCube experiment observed 28 high-energy neutrino events with $E_{\text{EM}} > 30$ TeV (with E_{EM} being the deposited electromagnetic-equivalent energy in detector), which is substantially larger than the number of events expected from the atmospheric backgrounds (which is $10.6^{+5.0}_{-3.6}$ events). Notably, the IceCube experiment found two events with deposit energy of \sim PeV (1041^{+143}_{-144} TeV and 1141^{+143}_{-133} TeV, nicknamed as Ernie and Bert, respectively), while no event with larger E_{EM} has been observed.

IceCube claims the existence of a source of high energy cosmic-ray neutrinos other than the atmospheric one [1]. We may consider particle-physics or astrophysical origin of such high energy neutrinos, the former of which is the subject of our study. (Possible astrophysical origins include Active Galactic Nuclei [2–4], γ -ray burst [5–8], hypernova remnants [9,10], star-forming galaxies [11], Galactic cosmic-rays [12–14], neutron-star mergers [15] and cosmogenic neutrinos [16,17]. For other astrophysical discussion, see also [18,19].) Although it is premature to make any conclusion, the negative observation of the events with larger energy deposit

may indicate that the cosmic-ray neutrino spectrum has a cutoff at the energy around \sim PeV. In fact, we also note that no event is observed in the energy bins between 0.4–0.63 PeV and 0.63–1 PeV. It may be a consequence of a peak of the cosmic-ray electron neutrino spectrum at \sim PeV; this is because, within experimental uncertainties, the deposited energy is equal to the energy of the initial-state neutrino for ν_e charged current events, while it is below the energy of the neutrino for other types of events.

From particle-physics point of view, the structure in the neutrino spectrum mentioned above may be realized with a new physics at the energy scale higher than \sim PeV. We consider such a case in this Letter. In particular, we study whether the decay of a new particle X with its mass $m_X \gtrsim 1$ PeV is responsible for the high-energy neutrino events observed by IceCube. For example, if neutrinos are produced by the decay of a particle X with $m_X \sim$ PeV in the present universe, a neutrino spectrum with the cutoff at \sim PeV may be obtained. The case where the dark-matter particle plays the role of X was considered in [20–22]. (For early discussion about related issues, see [23–29].) On the contrary, even if the mass of X is much larger than \sim PeV, there still exists a possibility to produce present cosmic-ray neutrinos with $E \sim$ PeV. This is because, if the decay of X occurs earlier than the present epoch, the energy of the neutrino produced by the X decay is redshifted.

In this Letter, we study cosmic-ray neutrinos produced by a long-lived particle X . We show that some of the high energy neutrino events observed by IceCube can be due to neutrinos produced by the decay of X . In particular, we discuss that the peak in the cosmic-ray neutrino spectrum may show up at $E \sim 1$ PeV even with the mass of X much larger than PeV, if the lifetime of X is shorter than the present age of the universe. For such a scenario

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to work, the decay of X is required to occur at $z \lesssim O(10^3)$ (with z being the redshift), which implies that the mass of X can be as large as 10^{10} GeV. We also discuss how the neutrino spectrum depends on the properties of X .

We start our discussion without assuming any particular model for the heavy particle X . Instead, we parametrize the properties of X with the following three quantities:

$$m_X, \quad \tau_X, \quad Y_X, \quad (1)$$

where m_X and τ_X are the mass and the lifetime of X , respectively. In addition, Y_X is the so-called yield variable

$$Y_X \equiv \left[\frac{n_X(t)}{s(t)} \right]_{t \ll \tau_X}, \quad (2)$$

with s being the entropy density; with Y_X , the number density of X is given by

$$n_X(t) = Y_X s(t) e^{-t/\tau_X}. \quad (3)$$

In the following, to make our point clearer, we concentrate on the case where the neutrino produced by the decay of X is monochromatic (with the energy \bar{E}_ν); the energy distribution of the neutrino produced by the decay of X is expressed as

$$\frac{dN_\nu^{(X)}}{dE} = \bar{N}_\nu \delta(E - \bar{E}_\nu), \quad (4)$$

where \bar{N}_ν is the number of the neutrino produced by the decay of one X . (For our numerical study, we take $\bar{E}_\nu = m_X/2$ and $\bar{N}_\nu = 1$.) Because we are interested in electron neutrino, with which the peak in the IceCube result may be explained, we assume that the decay of X produces a sizable amount of ν_e (after taking account of the effects of neutrino oscillation). With the monochromatic distribution, we will see that the peak in the cosmic-ray neutrino flux can be obtained.

Now, we discuss the flux of the cosmic-ray neutrinos produced by the decay of X . In the total flux, there exist two contributions:

$$\Phi_\nu(t, E) = \Phi_\nu^{(\text{Cosmo})}(t, E) + \Phi_\nu^{(\text{Galaxy})}(t, E), \quad (5)$$

where $\Phi_\nu^{(\text{Cosmo})}$ and $\Phi_\nu^{(\text{Galaxy})}$ are contributions from cosmological distance and from our Galaxy, respectively. (The neutrino number density is given by $n_\nu(t) = \int dE \Phi_\nu(t, E)$.)

The flux of high energy neutrinos from the cosmological distance obeys the following Boltzmann equation:

$$\frac{\partial \Phi_\nu^{(\text{Cosmo})}}{\partial t} = -2H \Phi_\nu^{(\text{Cosmo})} + H E \frac{\partial \Phi_\nu^{(\text{Cosmo})}}{\partial E} + S_\nu(t, E) - \gamma_\nu(t, E) \Phi_\nu^{(\text{Cosmo})}, \quad (6)$$

where H is the expansion rate of the universe, $S_\nu(t, E)$ is the source term:

$$S_\nu(t, E) = \frac{1}{4\pi} \frac{n_X(t)}{\tau_X} \frac{dN_\nu^{(X)}}{dE}, \quad (7)$$

with n_X being the number density of X , and $\gamma_\nu(t, E)$ is the scattering rate:

$$\gamma_\nu(t, E) = \frac{1}{16\pi^2 E^2} \int_0^\infty dk f_{\text{BG}}(k) \int_0^{4kE} ds \sigma_{\text{tot}}(s), \quad (8)$$

with f_{BG} being the distribution function of the background (target) particle and $\sigma_{\text{tot}}(s)$ the total scattering cross section of high energy neutrinos for the processes with the center-of-mass energy \sqrt{s} .

By solving the Boltzmann equation, the neutrino spectrum at the present epoch t_0 is given by

$$\Phi_\nu^{(\text{Cosmo})}(t_0, E) = \int_{-\infty}^{t_0} dt \left(\frac{a_0}{a(t)} \right)^{-2} D_\nu(E; z(t)) S_\nu(t, a_0 E/a(t)), \quad (9)$$

where a is the scale factor and $a_0 \equiv a(t_0)$. In addition,

$$D_\nu(E; z(t)) = \exp \left[- \int_t^{t_0} dt' \gamma_\nu(t', (1+z(t'))E) \right], \quad (10)$$

with

$$1+z(t) \equiv \frac{a_0}{a(t)}. \quad (11)$$

(Notice that the arguments of $D_\nu(E; z)$ are chosen to be the present neutrino energy E and the redshift at the time of the neutrino production z .) When the decay process produces monochromatic neutrino, we obtain

$$\Phi_\nu^{(\text{Cosmo})}(t_0, E) = \frac{1}{4\pi} \frac{\bar{N}_\nu Y_X s(t_0)}{\tau_X E} \left[\frac{e^{-\bar{t}/\tau_X} D_\nu(E; z(\bar{t}))}{H(\bar{t})} \right]_{1+z(\bar{t})=\bar{E}_\nu/E}. \quad (12)$$

In our analysis, the effects of secondary neutrinos produced by the scattering processes are not included. This can be justified as far as the damping factor $D_\nu(E; z)$ is close to 1. Hereafter, we mostly consider such a case.

In the case of our interest, important scattering processes of the high energy neutrinos are with background neutrinos, and hence $f_{\text{BG}}(k) = (e^{k/T_\nu} + 1)^{-1}$, where T_ν is the neutrino temperature and is related to the temperature of the background radiation as $T_\nu = (4/11)^{1/3} T$. Here we neglect the neutrino masses in the distribution function, since, as we will see later, the damping factor becomes effective only for $z \gtrsim O(10^3)$, where $T_\nu \gtrsim O(0.1 \text{ eV})$. Then, the scattering rate is given by

$$\gamma_\nu(t, E) = \frac{T_\nu(t)}{\pi^2} \int_0^\infty dk k \log(1 + e^{-k/T_\nu(t)}) \sigma_{\text{tot}}(s = 4kE). \quad (13)$$

For the calculation of the damping rate of ν_ℓ , σ_{tot} is obtained by taking account of the effects of the following processes:

- $\nu_\ell + \bar{\nu}_{\ell, \text{BG}} \rightarrow \nu_\ell + \bar{\nu}_\ell$,
- $\nu_\ell + \bar{\nu}_{\ell, \text{BG}} \rightarrow \ell + \bar{\ell}$,
- $\nu_\ell + \bar{\nu}_{\ell, \text{BG}} \rightarrow f + \bar{f}$, with $f \neq \nu_\ell, \ell$,
- $\nu_\ell + \bar{\nu}_{\ell', \text{BG}} \rightarrow \nu_\ell + \bar{\nu}_{\ell'}$, with $\ell \neq \ell'$,
- $\nu_\ell + \bar{\nu}_{\ell', \text{BG}} \rightarrow \ell + \bar{\ell}'$, with $\ell \neq \ell'$,
- $\nu_\ell + \nu_{\ell, \text{BG}} \rightarrow \nu_\ell + \nu_\ell$,
- $\nu_\ell + \nu_{\ell', \text{BG}} \rightarrow \nu_\ell + \nu_{\ell'}$, with $\ell \neq \ell'$,

where f denotes quarks and leptons, ℓ denotes charged leptons, and the subscript “BG” is for background neutrinos. (Cross sections for these processes are given in [Appendix A](#).) As we will see, for a cosmic-ray neutrino whose present energy is $E \sim 1 \text{ PeV}$, the damping becomes important only for $z \gtrsim O(10^3)$, for which the typical center-of-mass energy of the scattering process is $\sqrt{s} \gtrsim 1 \text{ GeV}$. Therefore, in our calculation of σ_{tot} , we neglect the masses of leptons as well as those of first- and second-generation quarks. We have checked that the damping factor $D_\nu(E; z)$ is insensitive to the behavior of $\sigma_{\text{tot}}(s)$ with $\sqrt{s} \lesssim$ a few GeV.

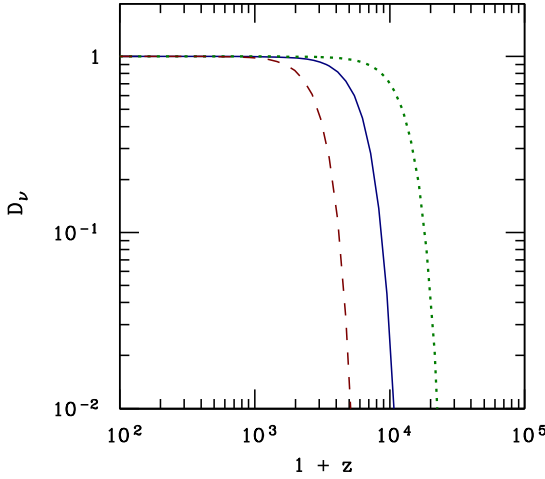


Fig. 1. The damping factor $D_v(E; z)$ defined in Eq. (10) as a function of the redshift $1+z$. The present energy of the neutrino is $E = 0.1$ (green, dotted), 1 (blue, solid), and 10 PeV (red, dashed) from right to left.

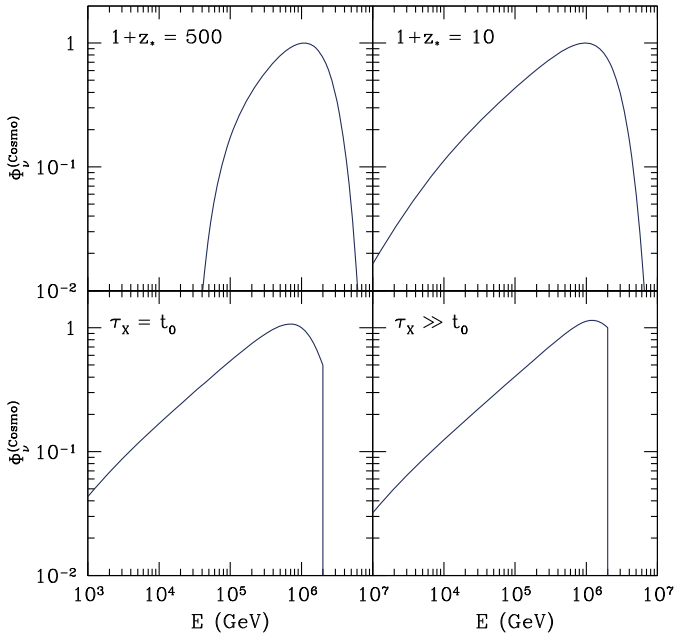


Fig. 2. Present cosmic-ray neutrino flux $\phi_v^{(\text{Cosmo})}$ given in Eq. (12) as a function of the present energy E . Here, we take $(\tau_X, \bar{E}_v) = (4.2 \times 10^{13} \text{ s}, 1000 \text{ PeV})$ (top-left, which corresponds to $1+z_* = 500$), $(1.7 \times 10^{16} \text{ s}, 20 \text{ PeV})$ (top-right, which corresponds to $1+z_* = 10$), $(t_0, 2 \text{ PeV})$ (bottom-left), and $(10^{29} \text{ s}, 2 \text{ PeV})$ (bottom-right). The normalization is arbitrary.

In Fig. 1, we plot the damping factor defined in Eq. (10) as a function of the redshift. For the present neutrino energy of $E \sim 1$ PeV, $D_v(E; z)$ is significantly suppressed for $z \gtrsim 0(10^3)$. For example, $D_v(E = 1 \text{ PeV}; z) < 0.5$ for $z \gtrsim 6 \times 10^3$.

In Fig. 2, we plot $\phi_v^{(\text{Cosmo})}(E)$ given in Eq. (12) for several values of τ_X . (Hereafter, we concentrate on the flux at present, so we omit t_0 from the argument.) The normalization of the spectrum in Fig. 2 is arbitrary. One can easily see a peak in the spectrum; the position of the peak depends on the initial energy of the neutrino as well as on the lifetime of X . If $\tau_X \ll t_0$ and $z_* \lesssim 10^3$ (for which the damping factor is negligible), the position of the peak is approximately given by

$$E_{\text{peak}}^{(\text{Cosmo})} \simeq 0.5 \times \frac{\bar{E}_v}{1+z_*} : \quad \text{for } \tau_X \ll t_0, \quad (14)$$

where

$$z_* = z(\tau_X). \quad (15)$$

This behavior can easily be understood by using the fact that most of X decays at the cosmic time of $t \sim \tau_X$, and that the energy of the emitted neutrino is redshifted by the factor of $\sim (1+z_*)^{-1}$. One can also see that, compared to the flux for $1+z_* = 10$, that for $1+z_* = 500$ shows a significant suppression for $E \lesssim 0.1$ PeV. This is due to the fact that the damping factor D_v becomes extremely small for $z \gtrsim 0(10^3)$. (The flux for $1+z_* = 10$ is almost unaffected by the damping.) For $1+z_* = 500$, $D_v \simeq 0.5$ for $E = 0.08$ PeV. For smaller E , secondary neutrinos produced by the scattering processes may also contribute and the flux may be affected. However, for such a redshift, the damping factor becomes close to 1 for $E \gtrsim 0.1$ PeV, so we believe that Eq. (12) well describes the contribution from cosmological distance in the energy region of our interest. Contrary to the case of $\tau_X \ll t_0$, the shape of the spectrum for $\tau_X \gg t_0$ is insensitive to τ_X as far as \bar{E}_v is fixed. We can also see that the shape of the spectrum for $\tau_X \gg t_0$ is quite different from that for $\tau_X \ll t_0$. In particular, $\phi_v^{(\text{Cosmo})}$ has a sharp edge at $E = \bar{E}_v$.

Next, we consider the contribution from the Milky-Way Galaxy. Because the Galactic contribution is not isotropic, we define

$$\tilde{\Phi}_v^{(\text{Galaxy})}(E, \hat{l}) = \frac{1}{4\pi} \frac{\bar{N}_v}{\tau_X} \delta(E - \bar{E}_v) \int_{\text{l.o.s.}} d\vec{l} n_X(\vec{l}), \quad (16)$$

where l.o.s. stands for the line-of-sight integral, and \hat{l} denotes the direction of the line-of-sight. We assume that the density of X is proportional to that of dark matter:

$$n_X(\vec{l}) = \frac{1}{m_X} \frac{\Omega_X}{\Omega_{\text{DM}}} \rho_{\text{DM}}(\vec{l}), \quad (17)$$

where Ω_X and Ω_{DM} are the density parameters of X and dark matter, respectively. In addition, $\rho_{\text{DM}}(\vec{l})$ is the energy density of dark matter in the Galaxy; in our numerical calculation, we adopt the NFW density profile [30,31]:

$$\rho_{\text{DM}}(r) = \rho_\odot \frac{r_\odot(r_c + r_\odot)^2}{r(r_c + r)^2}, \quad (18)$$

where r is the distance from the galactic center, $\rho_\odot \simeq 0.4 \text{ GeV/cm}^3$ is the local halo density, $r_c \simeq 20 \text{ kpc}$ is the core radius, and $r_\odot \simeq 8.5 \text{ kpc}$ is the distance between the galactic center and the solar system. Then, we define $\phi_v^{(\text{Galaxy})}$ as the directional average of $\tilde{\Phi}_v^{(\text{Galaxy})}$:

$$\phi_v^{(\text{Galaxy})}(E) \equiv \frac{1}{4\pi} \int d\Omega_i \tilde{\Phi}_v^{(\text{Galaxy})}(E, \hat{l}). \quad (19)$$

Now we are at the position to compare the predicted neutrino spectrum with the IceCube result. In our analysis, we particularly pay attention to the fact that IceCube has observed two events in the bin of $1 < E_{\text{EM}} < 1.6 \text{ PeV}$ while no event in $0.4 < E_{\text{EM}} < 0.63 \text{ PeV}$, $0.63 < E_{\text{EM}} < 1 \text{ PeV}$, and $E_{\text{EM}} > 1.6 \text{ PeV}$. Thus, we concentrate on the case where the spectrum of the neutrinos from X decay, which is assumed to contain a sizable fraction of ν_e , becomes largest at $\sim \text{PeV}$. To make a comparison with the IceCube result, we define the “averaged” neutrino spectrum for $1 < E < 1.6 \text{ PeV}$:

$$\bar{\Phi}_{v, \text{PeV}} \equiv \frac{1}{0.6 \text{ PeV}} \int_{1 \text{ PeV}}^{1.6 \text{ PeV}} dE \Phi_v(E), \quad (20)$$

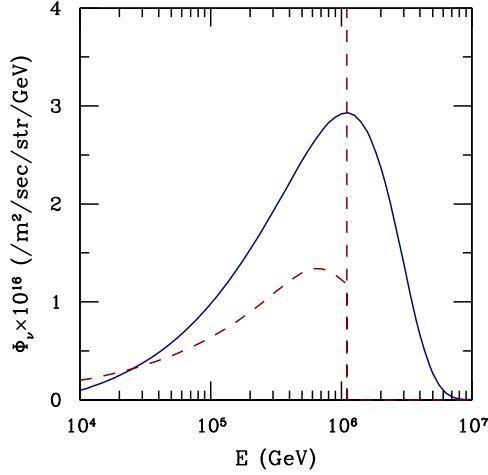


Fig. 3. The total present cosmic-ray neutrino flux from the long-lived particle X , normalized by $10^{-16} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV}^{-1}$. Here, we take $(\tau_X, \bar{E}_\nu, Y_X) = (5.2 \times 10^{14} \text{ s}, 2.2 \times 10^2 \text{ PeV}, 1.2 \times 10^{-26})$ (blue, solid, corresponding to $1 + z_* = 100$), and $(10^{29} \text{ s}, 1.1 \text{ PeV}, 4.0 \times 10^{-16})$ (red, dashed). The vertical line at $E = 1.1 \text{ PeV}$ is the contribution from the Galaxy.

and similar quantities $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Cosmo})}$ and $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}$ from $\phi_\nu^{(\text{Cosmo})}$ and $\phi_\nu^{(\text{Galaxy})}$, respectively. For $\tau_X \ll t_0$, $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}$ is negligibly small. On the contrary, for $\tau_X \gg t_0$, $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}$ becomes more important than $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Cosmo})}$ (as far as $1 < \bar{E}_\nu < 1.6 \text{ PeV}$). For $\bar{E}_\nu = 1.1\text{--}1.6 \text{ PeV}$, $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Cosmo})}/\bar{\phi}_{\nu, \text{PeV}}^{(\text{Galaxy})} \simeq 0.07\text{--}0.3$ when $\tau_X \gg t_0$. Using the fact that IceCube observed 2 events in the bin of $1 < E_{\text{EM}} < 1.6 \text{ PeV}$ within the live time of 662 days, we estimate

$$\bar{\phi}_{\nu, \text{PeV}}^{(\text{IceCube})} \simeq 3 \times 10^{-16} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV}^{-1}, \quad (21)$$

where we assumed that the total energy deposit is (almost) equal to the energy of the initial-state neutrino, which is the case of the charged current events of ν_e . Here, we used the effective area of 15 m^2 [1] as a reference value, although this value also includes the effects of neutral current. We take this flux as a canonical value for our study.

As we have mentioned, the neutrino flux is proportional to the yield variable Y_X . We estimate Y_X which gives $\bar{\phi}_{\nu, \text{PeV}} = 3 \times 10^{-16} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV}^{-1}$. For $\tau_X \ll t_0$, for which $\bar{\phi}_{\nu, \text{PeV}}^{(\text{Cosmo})}$ is negligible, we choose m_X so that $E_{\text{peak}}^{(\text{Cosmo})} = 1.1 \text{ PeV}$ (which is close to the deposited energies of the most energetic events, i.e., Ernie and Bert). Then, the best-fit value of Y_X to realize Eq. (21) is given by

$$Y_X \simeq 1 \times 10^{-26} \times \bar{N}_\nu^{-1} : \quad \tau_X \ll t_0. \quad (22)$$

Notice that, for $\tau_X \ll t_0$, the best-fit value is insensitive to the lifetime of X . On the contrary, for $\tau_X \gg t_0$, we choose $\bar{E}_\nu = 1.1 \text{ PeV}$ and obtain

$$Y_X \simeq 4 \times 10^{-16} \times \bar{N}_\nu^{-1} \left(\frac{\tau_X}{10^{29} \text{ s}} \right) : \quad \tau_X \gg t_0. \quad (23)$$

In the case of $\tau_X \gg t_0$, Y_X corresponds to the present yield value of X . Then, the present mass density of X is estimated as $\Omega_X \simeq 6 \times 10^{14} \times Y_X (m_X/1 \text{ PeV})$. Combining this relation with Eq. (23), we can see that dark matter may play the role of X if $\tau_X \sim O(10^{29} \text{ s})$ [20–22]. On the contrary, scenarios with longer lifetime do not work because of the over-closure of the universe.

In Fig. 3, we show the neutrino spectrum for $(\tau_X, \bar{E}_\nu) = (5.2 \times 10^{14} \text{ s}, 2.2 \times 10^2 \text{ PeV})$ and $(10^{29} \text{ s}, 1.1 \text{ PeV})$; the yield variable is determined so that $\bar{\phi}_{\nu, \text{PeV}}$ is equal to $\bar{\phi}_{\nu, \text{PeV}}^{(\text{IceCube})}$ given in Eq. (21).

We can see that the enhancement of the flux at $E \sim 1 \text{ PeV}$ is possible both for $\tau_X \ll t_0$ and $\tau_X \gg t_0$.

One of the important check points of the present scenario, in particular for the case of $\tau_X \ll t_0$, is the flux at the “Glashow resonance” [32]. For $E \simeq 6.3 \text{ PeV}$, the event rate of IceCube for ν_e is enhanced because the center-of-mass energy hits the W -boson pole [33,34]. Indeed, the effective area of the IceCube experiment for $E \simeq 6.3 \text{ PeV}$ is about 40 times larger than that for $E \simeq 1 \text{ PeV}$. One might wonder if the present scenario is consistent with the negative observation of the events for such an energy region because, if $\tau_X \ll t_0$, the spectrum is non-vanishing even at $E \sim 1 \text{ PeV}$. Notably, the flux for $E \gg E_{\text{peak}}^{(\text{Cosmo})}$ is exponentially suppressed because neutrinos with $E \gg E_{\text{peak}}^{(\text{Cosmo})}$ are produced when $t \gg \tau_X$. In fact, the flux at $E \simeq 6.3 \text{ PeV}$ is quite sensitive to the value of z_* . With the position of the peak being fixed as $E_{\text{peak}}^{(\text{Cosmo})} \simeq 1 \text{ PeV}$, the ratio $\phi_\nu^{(\text{Cosmo})}(6.3 \text{ PeV})/\phi_\nu^{(\text{Cosmo})}(E_{\text{peak}}^{(\text{Cosmo})})$ becomes smaller for larger value of z_* ; this can be understood from the fact that $\phi_\nu^{(\text{Cosmo})}(E)$ is proportional to e^{-t_E/τ_X} , where t_E is the time satisfying $E = (1 + z(t_E))^{-1} \bar{E}_\nu$ (see Eq. (12)). For the case where X decays in radiation- and matter-dominated epochs, for example, this quantity is given by $e^{-t_E/\tau_X} = e^{-(E/E_*)^2}$ and $e^{-(E/E_*)^{3/2}}$, respectively, with $E_* \equiv (1 + z_*)^{-1} \bar{E}_\nu$. For $E_{\text{peak}}^{(\text{Cosmo})} = 1 \text{ PeV}$ (1.1 PeV), $\phi_\nu^{(\text{Cosmo})}(6.3 \text{ PeV})/\phi_\nu^{(\text{Cosmo})}(E_{\text{peak}}^{(\text{Cosmo})})$ is given by 0.014, 0.012, and 0.001, (0.027, 0.024, and 0.003) for $1 + z_* = 10, 100$, and 1000, respectively. Using the fact that IceCube has not observed any event at around the Glashow resonance, relatively large value of z_* may be preferred. However, the statistics are still poor, and it is premature to exclude the possibility of small z_* . With more data, IceCube may see events at around $E \simeq 6.3 \text{ PeV}$ in particular in the case with small z_* .

In our discussion, we have concentrated on the case where the high energy neutrino spectrum is peaked at $\sim \text{PeV}$. Interpretation of the IceCube events with E_{EM} smaller than $\sim \text{PeV}$ may be an issue. For those events, it should be noted that the deposited energy E_{EM} is in general not equal to the energy of the initial-state neutrino; this is because some amount of the energy may be carried away by neutrinos in the final state. Thus, even with the present setup, events with $E_{\text{EM}} < 1 \text{ PeV}$ can be induced by charged current interaction of neutrinos other than ν_e and neutral current ones, as well as by the neutrinos with E smaller than PeV . Unfortunately, the expected distribution of E_{EM} for the present neutrino spectrum is not available. Although it is beyond the scope of this Letter, such a study should be given elsewhere to test the present scenario. For the excess of the events with $E_{\text{EM}} < 1 \text{ PeV}$, we also point out that we may also consider non-monochromatic initial neutrino injection as another possibility [21,22].

Here, let us comment on the possibilities to acquire information about the lifetime of X in the present scenario. Because the detailed shape of the spectrum depends on the lifetime, it may be possible to distinguish the cases with $\tau_X \ll t_0$ and $\tau_X \gg t_0$. As we can see, the spectrum smoothly continues to $E > E_{\text{peak}}$ if $\tau_X \ll t_0$. On the contrary, for $\tau_X \gg t_0$, the neutrino flux is dominated by the one originating in the Galaxy. Then, the spectrum is sharply peaked at $E = \bar{E}_\nu$. Thus, if the spectrum of the neutrinos is precisely determined in the future, it will provide important information about the lifetime of X . Another possibility is to use the directional information about the high energy neutrinos. In the case of $\tau_X \ll t_0$, high energy neutrinos originate in the decay of X at high redshift so that they are isotropic. On the contrary, for the case of $\tau_X \gg t_0$, a large fraction of the high energy neutrino events are Galactic origin. Consequently, the neutrino flux is enhanced for the direction of the Galactic center. We define θ as the angle between the direction of the high-energy neutrino and that of the

Galactic center, with $\theta = 0$ being the direction of the Galactic center. Then, we calculate $\bar{\Phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}(\theta < 90^\circ)$ and $\bar{\Phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}(\theta > 90^\circ)$, which are angular-averaged neutrino fluxes in the regions of $\theta < 90^\circ$ and $\theta > 90^\circ$, respectively, for $1 < E < 1.6$ PeV (see Eq. (20)). For $\tau_X \gg t_0$, $\bar{\Phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}(\theta < 90^\circ)/\bar{\Phi}_{\nu, \text{PeV}}^{(\text{Galaxy})}(\theta > 90^\circ) \simeq 2$. Thus, significant angular dependence is expected for $\tau_X \gg t_0$, while the flux is isotropic for $\tau_X \ll t_0$. Experimental determination of the directional distribution is important to distinguish the scenarios with $\tau_X \ll t_0$ and $\tau_X \gg t_0$.

Before closing this Letter, several comments are in order. First, we consider a possible scenario to produce X in the early universe. As we have seen, for $\bar{\Phi}_{\nu, \text{PeV}} = 3 \times 10^{-16} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV}^{-1}$, very small value of Y_X is required. In particular, if $\tau_X \ll t_0$, $Y_X \sim O(10^{-(26-27)})$ is necessary, which is much smaller than the typical thermal relic abundance. Even if we assume that X is in the thermal bath, however, Y_X can be significantly suppressed with large entropy production. A mini inflation after the first inflation (which is responsible for the density perturbation of the universe) may be an example. Another possibility is to introduce a field which has a very small branching ratio into X ; if such a field is produced in the early universe, X can be produced by its decay.

We also comment on possible constraints from high energy cosmic rays, in particular, γ -ray. With the mass of X as large as $\sim 10^{10}$ GeV, electroweak gauge bosons and charged leptons are also produced by electroweak jet cascade even if X dominantly decays into neutrinos [35]. (If X has decay modes into gauge bosons or into charged leptons, they also contribute.) If too much γ -ray is generated, the present scenario conflicts with the observations of extragalactic γ -ray. In the present scenario, however, we expect that the flux is small enough by estimating the total amount of energy injection due to the decay of X . Assuming that the productions of electroweak gauge bosons and charged leptons are subdominant compared to the neutrino production, the energy density of radiation (i.e., γ -ray) from X should be smaller than $E_\gamma^2 \Phi_\gamma \lesssim O(10^{-4} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV})$ (see Eq. (21)). Notice that, if the production process of γ -ray is suppressed, the flux becomes smaller; this is the case for the electroweak jet cascade processes. For the energy range for which the measurement of the extragalactic γ -ray flux is available (i.e., $E_\gamma \lesssim 100$ GeV), the γ -ray flux in the present scenario is found to be smaller than the observed one (which is $E_\gamma^2 \Phi_\gamma \gtrsim O(10^{-3} \text{ m}^{-2} \text{ s}^{-1} \text{ str}^{-1} \text{ GeV})$) [36]. Thus, we believe that the present scenario is not excluded by the current measurements of high energy extragalactic γ -ray flux. With future improvements of the measurements, signals of the decay of X may be seen. The detailed understanding of the signals requires a precise calculation of the spectrum of γ -ray, which is beyond the scope of this Letter.

In summary, in this Letter, we have studied the cosmic-ray neutrinos produced by a long-lived particle X . The PeV neutrino events observed by IceCube may be due to the neutrinos produced by a heavy particle X . We have discussed that such a scenario works even with the mass of X much higher than $\sim \text{PeV}$ if the lifetime is shorter than the present cosmic time. To make such a scenario viable, we have seen that the decay of X should occur at $z \lesssim O(10^3)$, and that the scale of the new physics responsible for the IceCube events can be as large as $O(10^{9-10} \text{ GeV})$. Detailed study of the propagation of neutrinos taking into account the effects of the secondary neutrinos is necessary to understand the case with $m_X \gtrsim O(10^{9-10} \text{ GeV})$.

Finally, we discuss particle-physics models which contain a long-lived particle decaying into neutrino. Even assuming that X is a neutral scalar particle, one may consider the case where X is embedded into $SU(2)_L$ triplet (with the hypercharge of $+1$) [20]. Another possibility is to introduce $SU(2)_L$ doublet boson (with the hypercharge of $-1/2$) other than ordinary Higgs boson, which

couples to lepton doublet and right-handed neutrino ν_R^c . Then, identifying the neutral component of the doublet as X , the decay process $X \rightarrow \nu_L \nu_R^c$ becomes possible if the neutrino mass is Dirac type or the Majorana mass of ν_R^c is small enough. So far, we have considered the case where X is neutral. For the case of $\tau_X \ll t_0$, however, X may be charged (or even colored) because the constraints on stable superheavy charged particle (in particular, those using sea water [37]) do not apply. From particle-physics point of view, there exist various beyond-standard-model physics which contain the candidate of X , in particular when the scale of the new physics is around $O(10^{9-10} \text{ GeV})$. One example is a fermion in Peccei–Quinn (PQ) sector [38,39] (or its supersymmetric partner) in hadronic axion model [40,41]. The PQ (s)fermion may be stable if it has no mixing with standard-model fermions. With introducing a very small mixing, the PQ (s)fermion decays into standard-model particles with a very long lifetime. Similar argument holds for a messenger (s)fermion in gauge-mediation supersymmetry breaking model [42–44]. Some of those particles may play the role of X if they decay into neutrinos at $z \lesssim O(10^3)$.

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Appendix A

In this appendix, we summarize the cross sections for the neutrino scattering processes which are relevant for the calculation of the damping rate of high energy neutrinos. The fermion masses (in particular, lepton masses) are neglected except in Eq. (26).

- $\nu_\ell + \bar{\nu}_{\ell, \text{BG}} \rightarrow \nu_\ell + \bar{\nu}_\ell$:

$$\begin{aligned} \sigma = & \frac{g_Z^4}{192\pi} \frac{s}{(s - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} \\ & + \frac{g_Z^4}{64\pi s} [x_Z^{-1} + 2 - 2(1 + x_Z) \log(1 + x_Z^{-1})] \\ & + \frac{g_Z^4}{64\pi} \frac{s}{(s - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} (1 - x_Z) \\ & \times [3 + 2x_Z - 2(1 + x_Z)^2 \log(1 + x_Z^{-1})], \end{aligned} \quad (24)$$

where, here and hereafter, $x_V \equiv m_V^2/s$ (with $V = W$ and Z).

- $\nu_\ell + \bar{\nu}_{\ell, \text{BG}} \rightarrow \ell + \bar{\ell}$:

$$\begin{aligned} \sigma = & \frac{g_Z^2 (g_{Z, \ell_L}^2 + g_{Z, \ell_R}^2)}{48\pi} \frac{s}{(s - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} \\ & + \frac{g_Z^4}{16\pi s} [x_W^{-1} + 2 - 2(1 + x_W) \log(1 + x_W^{-1})] \\ & + \frac{g_Z g_{Z, \ell_L} g_{Z, \ell_R}^2}{16\pi} \frac{s}{(s - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} (1 - x_Z) \\ & \times [3 + 2x_W - 2(1 + x_W)^2 \log(1 + x_W^{-1})]. \end{aligned} \quad (25)$$

- $\nu_\ell + \bar{\nu}_{\ell,\text{BG}} \rightarrow f + \bar{f}$ ($f \neq \nu_\ell, \ell$):

$$\sigma = \frac{g_Z^2}{48\pi} \frac{(g_{Z,f_L}^2 + g_{Z,f_R}^2)(s - m_f^2) + 6g_{Z,f_L}g_{Z,f_R}m_f^2}{(s - m_Z^2)^2 + m_Z^2\Gamma_Z^2} \times \sqrt{1 - \frac{4m_f^2}{s}}, \quad (26)$$

where m_f is the mass of f .

- $\nu_\ell + \bar{\nu}_{\ell',\text{BG}} \rightarrow \nu_\ell + \bar{\nu}_{\ell'} (\ell \neq \ell')$:

$$\sigma = \frac{g_Z^4}{64\pi s} [x_Z^{-1} + 2 - 2(1 + x_Z) \log(1 + x_Z^{-1})]. \quad (27)$$

- $\nu_\ell + \bar{\nu}_{\ell',\text{BG}} \rightarrow \ell + \bar{\ell}' (\ell \neq \ell')$:

$$\sigma = \frac{g_Z^4}{16\pi s} [x_W^{-1} + 2 - 2(1 + x_W) \log(1 + x_W^{-1})]. \quad (28)$$

- $\nu_\ell + \nu_{\ell,\text{BG}} \rightarrow \nu_\ell + \nu_\ell$:

$$\sigma = \frac{g_Z^4}{64\pi s} \frac{1}{x_Z(1 + x_Z)} + \frac{g_Z^4}{32\pi s} \frac{1}{1 + 2x_Z} \log(1 + x_Z^{-1}). \quad (29)$$

- $\nu_\ell + \nu_{\ell',\text{BG}} \rightarrow \nu_\ell + \nu_{\ell'} (\ell \neq \ell')$:

$$\sigma = \frac{g_Z^4}{64\pi s} \frac{1}{x_Z(1 + x_Z)}. \quad (30)$$

In the above expressions, g_Z is the gauge coupling constant for $SU(2)_L$, $g_Z \equiv \sqrt{g_2^2 + g_1^2}$ (with g_1 being the gauge coupling constant for $U(1)_Y$), and

$$g_{Z,u_L} = \frac{1}{2}g_Z - \frac{2}{3}\frac{g_1^2}{g_Z}, \quad (31)$$

$$g_{Z,u_R} = -\frac{2}{3}\frac{g_1^2}{g_Z}, \quad (32)$$

$$g_{Z,d_L} = -\frac{1}{2}g_Z + \frac{1}{3}\frac{g_1^2}{g_Z}, \quad (33)$$

$$g_{Z,d_R} = \frac{1}{3}\frac{g_1^2}{g_Z}, \quad (34)$$

$$g_{Z,\ell_L} = -\frac{1}{2}g_Z + \frac{g_1^2}{g_Z}, \quad (35)$$

$$g_{Z,\ell_R} = \frac{g_1^2}{g_Z}, \quad (36)$$

$$g_{Z,\nu_L} = \frac{1}{2}g_Z, \quad (37)$$

$$g_{Z,\nu_R} = 0. \quad (38)$$

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